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Phases in the MSSM, Electric Dipole Moments and Cosmological Dark Matter

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Abstract

We consider the effect of CP violating phases in the MSSM on the relic density of the lightest supersymmetric particle (LSP). In particular, we find that the upper limits on the LSP mass are relaxed when phases in the MSSM are allowed to take non-zero values when the LSP is predominantly a gaugino (bino). Previous limits of $m_{\tilde{B}} \lesssim 250$ GeV for $\Omega h^2 < 0.25$ can be relaxed to $m_{\tilde{B}} \lesssim 650$ GeV. We also consider the additional constraints imposed by the neutron and electron electric dipole moments induced by these phases. Though there is some restriction on the phases, the bino mass may still be as large as ~ 350 GeV and certain phases can be arbitrarily large.

It is well known that by considering the cosmological relic density of stable particles, one can establish mass limits on these particles. For example, light neutrinos ($m_\nu < 1$ MeV) contribute too much to the overall cosmological mass density if $m_\nu \gtrsim 25$ eV for $\Omega h^2 < 1/4$ [1]. The relic density of heavier neutrinos are determined by their annihilation cross section and yield lower bounds to neutrino masses $m_\nu \gtrsim 4\text{--}7$ GeV [2, 3]. In the minimal supersymmetric standard model (MSSM), because of the many unknown parameters it becomes a significantly more complicated task to set limits on the mass of the lightest particle, the LSP, e.g. the annihilation cross section will depend on several parameters which determine the identity of the LSP. Here we will consider only neutralinos as the LSP (for a recent discussion on the possibilities for sneutrino dark matter, see [4]). The mass limits will then depend sensitively on the masses of the scalars [5], and in general, limits are obtained on portions of the parameter space [6]. For example when parameters are chosen so that the LSP is a gaugino (when the supersymmetry breaking gaugino masses $M_{1,2}$ are taken to be smaller than the supersymmetric Higgs mixing mass, μ), for a given set of scalar masses, the requirement that $\Omega h^2 < 1/4$ places a lower bound on the gaugino mass. For small gaugino masses, this corresponds to a lower limit on the photino mass [5]. At larger gaugino masses (but still smaller than μ), the LSP is a bino [7]. As the bino mass is increased, there is an upper limit $m_{\tilde{B}} \lesssim 250$ GeV [7, 8]. In [9], it was shown that this upper limit is sensitive to the level of sfermion mixing. Here, we find that the upper limit to the bino mass is relaxed when phases in the MSSM, and in particular the phases associated with the off-diagonal sfermion masses, are allowed to take non-zero values. Now, the upper limit is increased to $m_{\tilde{B}} \lesssim 650$ GeV.

There has been a considerable amount of work concerning phases in the MSSM. For the most part these phases are ignored because they tend to induce large electric dipole moments for the neutron [11, 12]. To suppress the electric dipole moments, either large scalar masses (approaching 1 TeV) or small angles (of order 10^{-3} , when all SUSY masses are of order 100 GeV) are required. For the most part, the community has opted for the latter, though the possibilities for large phases was recently considered in [13]. To reconcile large phases with small electric dipole moments, some of the sparticle masses are required to be heavy. In [13], either large sfermion or neutralino masses (or both) were required. However, unless R -parity is broken and the LSP is not stable, one would require that the sfermions be heavier than the neutralinos. If they are much heavier, this would result in an excessive relic density of neutralinos. The object of this letter is to determine the relationship between potentially large phases in the MSSM and the relic density while remaining consistent with experimental bounds on the electric dipole moments.

In the MSSM, the possibility for new phases arises from a number of sources. First, in the superpotential, there is the Higgs mixing mass, μ . There are also several parameters associated with supersymmetry breaking: gaugino masses M_i , $i = 1-3$ (we will assume GUT conditions on these masses so that we need only consider one of them, M_2); and in the scalar sector, soft scalar masses; soft bilinear, $B\mu$, and trilinear, A_f , terms. Not all of these phases are physical [12]. It is common to rotate away the phase of the gaugino masses, and to make $B\mu$ real, which ensures that the vacuum expectation values of the Higgs fields are real. We will also, for simplicity, ignore generation mixing in the sfermion sector though we will include left-right mixing. If furthermore we assume that all of the A_f 's are equal, and simply label them by A , we are left with two independent phases, the phase of A , θ_A , and the phase of μ , θ_μ . The phase of B is fixed by the condition that $B\mu$ is real.

As both the relic cosmological density of neutralinos and the electric dipole moments are strongly dependent on the sfermion masses, it will be very useful to identify the combination of phases that enter the sfermion mass² matrix. We take the general form of the sfermion mass² matrix to be [14]

$$\begin{pmatrix} M_L^2 + m_f^2 + \cos 2\beta(T_{3f} - Q_f \sin^2 \theta_W)M_Z^2 & m_f \overline{m}_f e^{i\gamma_f} \\ m_f \overline{m}_f e^{-i\gamma_f} & M_R^2 + m_f^2 + \cos 2\beta Q_f \sin^2 \theta_W M_Z^2 \end{pmatrix} \quad (1)$$

where $M_{L(R)}$ are the soft supersymmetry breaking sfermion mass which we have assumed are generation independent and generation diagonal and hence real. We will also assume $M_L = M_R$. m_f is the mass of the fermion f , $\tan \beta$ is the ratio of Higgs vevs, and $R_f = \cot \beta$ ($\tan \beta$) for weak isospin $+1/2$ ($-1/2$) fermions. Due to our choice of phases, there is a non-trivial phase associated with the off-diagonal entries, which we denote by $m_f(\overline{m}_f e^{i\gamma_f})$, of the sfermion mass² matrix, and

$$\overline{m}_f e^{i\gamma_f} = R_f \mu + A^* = R_f |\mu| e^{i\theta_\mu} + |A| e^{-i\theta_A} \quad (2)$$

For a given value of μ and A , there are then two phases which can be distinguished (by R_f), and we denote them by γ_t and γ_b . The sfermion mass² matrix is diagonalized by the unitary matrix U ,

$$U = \begin{pmatrix} u & v \\ -v^* & u \end{pmatrix} \quad (3)$$

where $u^2 + |v|^2 = 1$, and we have taken u real.

Previously [9] we considered the effect of sfermion mixing on the relic density when the neutralino is mostly gaugino, and in particular a bino. The LSP is a bino whenever $M_2 < \mu$ and $M_2 \gtrsim 200$ GeV; for smaller M_2 , the bino is still the LSP for large enough values of

μ . The bino portion of the $M_2 - \mu$ LSP parameter plane is attractive, as it offers the largest possibility for a significant relic density [7, 8, 15]. The complementary portion of the parameter plane, with $\mu < M_2$, only gives a sizable density in a limited region, due to the large annihilation cross sections to W^+W^- and ZZ and due to co-annihilations [16] with the next lightest neutralino (also a Higgsino in this case), which is nearly degenerate with the LSP [17]. We will also focus on binos as the LSP here. In addition to resulting in a sizable relic density, the analysis is simplified by the fact that in the nearly pure bino region, the composition and mass of the LSP is not very sensitive to the new phases. However, as we will now show, the relic density, which is determined primarily by annihilations mediated by sfermion exchange, is quite sensitive to the phases, γ_t and γ_b .

We begin by exploring the effect of the new phases on the relic density of binos. In the absence of these phases and in the absence of sfermion mixing, there is an upper limit [7, 8] on the bino mass of $m_{\tilde{B}} \lesssim 250$ GeV for $\Omega h^2 < 1/4$. (This upper limit is somewhat dependent on the value of the top quark mass. In [7, 9], we found that the upper limit was ~ 250 GeV for $m_t \sim 100$ GeV. When $m_t = 174$ GeV, the upper limit is 260 GeV. For $m_t \sim 200$ GeV, this limit is increased to $m_{\tilde{B}} \lesssim 300$ GeV. Furthermore, one should be aware that there is an upward correction of about 15% when three-body final states are included [10] which raises the bino mass limit to about 350 GeV. This latter correction would apply to the limits discussed below though it has not been included.) As the bino mass is increased, the sfermion masses, which must be larger than $m_{\tilde{B}}$, are also increased, resulting in a smaller annihilation cross section and thus a higher relic density. At $m_{\tilde{B}} \simeq 250$ GeV, even when the sfermion masses are equal to the bino mass, $\Omega h^2 \sim 1/4$. Note that this provides an upper bound on the sfermion masses as well, since the mass of the lightest sfermion is equal to the mass of the bino, when the bino mass takes its maximum value. When sfermion mixing is included [9], the limits, which now depend on the magnitude of the off-diagonal elements $m_f \bar{m}_f$, are modified. We find that this upper limit is relaxed considerably when the phases are allowed to take non-zero values.

The dominant contribution to bino annihilation is due to sfermion exchange and is derived from the bino-fermion-sfermion interaction Lagrangian,

$$\begin{aligned} \mathcal{L}_{f\tilde{f}\tilde{B}} &= \frac{1}{\sqrt{2}}g' \left(Y_R \bar{f} P_L \tilde{B} \tilde{f}_R + Y_L \bar{f} P_R \tilde{B} \tilde{f}_L \right) + \text{h.c.} \\ &= \frac{1}{\sqrt{2}}g' \bar{f} \left(\tilde{f}_1 x P_L + \tilde{f}_2 w P_L + \tilde{f}_1 y P_R - \tilde{f}_2 z P_R \right) \tilde{B} + \text{h.c.} \end{aligned} \quad (4)$$

where $x = -Y_R v^*$, $y = Y_L u$, $w = Y_R u$, $z = -Y_L v$, $P_{R,L} = (1 \pm \gamma_5)/2$, $Y_R = 2Q_f$ and $Y_L = 2(Q_f - T_{3f})$ where Q_f is the fermion charge and T_{3f} is the fermion weak isospin.

We compute the relic density by using the method described in ref. [3]. We expand $\langle\sigma v_{rel}\rangle$ in a Taylor expansion in powers of $T/m_{\tilde{B}}$

$$\langle\sigma v_{rel}\rangle = a + b (T/m_{\tilde{B}}) + O\left((T/m_{\tilde{B}})^2\right) \quad (5)$$

The coefficients a and b are given by

$$a = \sum_f v_f \tilde{a}_f \quad (6)$$

$$b = \sum_f v_f \left[\tilde{b}_f + \left(-3 + \frac{3m_f^2}{4v_f^2 m_{\tilde{B}}^2} \right) \tilde{a}_f \right] \quad (7)$$

where \tilde{a}_f and \tilde{b}_f are computed from the expansion of the matrix element squared in powers of p , the incoming bino momentum, and $v_f = (1 - m_f^2/m_{\tilde{B}}^2)^{1/2}$ is a factor from the phase space integrals.

We summarize the result by quoting the computed expression for \tilde{a}_f :

$$\tilde{a}_f = \frac{g'^4}{128\pi} \left| \frac{\Delta_1(m_f w^2 - 2m_{\tilde{B}} w z + m_f z^2) + \Delta_2(m_f x^2 + 2m_{\tilde{B}} x y + m_f y^2)}{\Delta_1 \Delta_2} \right|^2 \quad (8)$$

where $\Delta_i = m_{\tilde{f}_i}^2 + m_{\tilde{B}}^2 - m_f^2$. The result for \tilde{b}_f is too lengthy for presentation here, but was computed and used in the numerical integrations to obtain the relic densities. The results reduce to the results quoted in [9] in the limit of zero phases.

We show our results in Figures 1 and 2 for the upper and lower limits on $m_{\tilde{B}}$ as a function of the magnitude of the off-diagonal term in the top-squark mass² matrix, \overline{m}_t , given the conditions 1) $\Omega_{\tilde{B}} h^2 < 1/4$, 2) the lightest sfermion is heavier than the bino, and 3) the lightest sfermion is heavier than 74 GeV[18]. In both figures we have taken $\tan\beta = 2$ and $m_{\text{top}} = 174$ GeV. In Figure 1, $|\mu| = 3000$ GeV and in Figure 2, $|\mu| = 1000$ GeV. The various curves are labeled by the value of γ_b assumed, and in addition, $m_{\tilde{B}}$ has been maximized for all allowed values of θ_μ . The lower limit on $m_{\tilde{B}}$ assumes $\gamma_b = \pi/2$. As one can see, when γ_b is allowed to take its maximal value of $\pi/2$, the upper limits are greatly relaxed to $m_{\tilde{B}} \lesssim 650$ GeV. With $|\mu|$ and γ_b fixed, for a given value of \overline{m}_t and θ_μ all of the remaining quantities such as $|A|$, θ_A , γ_t , and \overline{m}_b are determined (though not necessarily uniquely, as some are double valued).

The reason for the change in bounds becomes evident if we consider $\Delta_1 \approx \Delta_2 \equiv \Delta$ and $m_f \ll m_{\tilde{B}}$ in (8). Then

$$\tilde{a}_f \approx \frac{(g')^4}{8\pi\Delta^2} m_{\tilde{B}}^2 Y_L^2 Y_R^2 u^2 \text{Im}(v)^2 + O(m_f m_{\tilde{B}})$$

$$= \frac{(g')^4}{8\pi\Delta^2} m_B^2 Y_L^2 Y_R^2 u^2 |v|^2 \sin^2 \gamma_f + O(m_f m_{\tilde{B}}) \quad (9)$$

The size of $u|v|$ depends on the quantity $r \equiv m_f \bar{m}_f / (\cos 2\beta (T_3 - 2Q_f \sin^2 \theta_W) M_Z^2)$; for $r \approx 1$, $u|v| \approx .45$, while for $r \ll 1$, $u|v| \approx r$. Note that a non-zero value for $\sin \gamma_f$ removes the p-wave suppression for the fermion f (to a much greater extent than sfermion mixing alone [9]) and greatly enhances the annihilation cross-section. Of course, for annihilations through the top quark channel, the p-wave suppression is not terribly strong, as m_t is not much smaller than $m_{\tilde{B}}$.

For large \bar{m}_t , the diagonal mass terms M_L^2 must be taken large to ensure that the mass of the lightest stop is $\gtrsim m_{\tilde{B}}$. This drives up the masses of the other sfermions and suppresses their contribution to the annihilation. As \bar{m}_t is decreased, M_L^2 must drop, the other sfermions begin to contribute and the upper bound on $m_{\tilde{B}}$ is increased. In particular, annihilation to μ 's and τ 's becomes important, since

$$Y_L^2 Y_R^2 \Big|_{\mu, \tau} : Y_L^2 Y_R^2 \Big|_{c, t} : Y_L^2 Y_R^2 \Big|_{s, b} = 81 : 4 : 1 \quad (10)$$

Decreasing γ_b reduces the effect of μ 's and τ 's, and this can be seen as a decrease in the upper bounds in Figures 1 and 2. For \bar{m}_t sufficiently small, the stops become unmixed, diminishing somewhat their contribution and slightly decreasing the upper bound on $m_{\tilde{B}}$.

Lastly, we mention how the top contribution depends on γ_t . For the large values of $m_{\tilde{B}}$ allowed for $\bar{m}_t \lesssim 1500$ GeV, top annihilation feels an enhancement from a non-zero $\sin \gamma_t$ as in (9) (but note that for the top, the higher order terms in m_f are not small). However, except in a narrow region near $\bar{m}_t = 0$, the two stop eigenstates \tilde{t}_1 and \tilde{t}_2 have a large mass splitting, and the bino couples with different strengths to \tilde{t}_1 and \tilde{t}_2 . Taking γ_t away from π decreases the coupling to the lighter stop (which can offer a much greater contribution to the annihilation than can the heavier stop), and we find that larger annihilation rates into tops are generally found nearer to $\gamma_t = \pi$ than $\gamma_t = \pi/2$. In Figure 2, the sudden drop in the $\theta_\mu = \pi/8$ curve at 1500 GeV occurs because for $\bar{m}_t > (\tan \beta - \cot \beta) |\mu|$, $\gamma_b - \pi/2 < \gamma_t < \gamma_b + \pi/2$, and so for $\theta_\mu = \pi/8$, γ_t is prevented from approaching π .

We turn now to the calculation of the electric dipole moments of the neutron and the electron. The EDM's of the electron and quarks receive contributions from one-loop diagrams involving the exchange of sfermions and either neutralinos, charginos, or (for the quarks) gluinos. In the case of the neutron EDM, there are additional operators besides the quark electric dipole operator, $O_\gamma = \frac{1}{4} \bar{q} \sigma_{\mu\nu} q \tilde{F}^{\mu\nu}$ [11] which contribute. They are the gluonic operator $O_G = -\frac{1}{6} f^{abc} G_a G_b \tilde{G}_c$ [19] and the quark color dipole operator, $O_q = \frac{1}{4} \bar{q} \sigma_{\mu\nu} q T^a \tilde{G}_a^{\mu\nu}$ [20]. The gluonic operator is the smallest [21, 20] when all mass scales are taken to be equal.

These three operators are conveniently compared to one another in [22] and relative to the gluino exchange contribution to the O_γ operator, it is found that $O_\gamma : O_q : O_g = 21 : 4.5 : 1$.

Because of the reduced importance of the additional operators contributing to the neutron EDM, we will only include the three contributions to the quark electric dipole moment. The necessary CP violation in these contributions comes from either γ_f in the sfermion mass matrices or θ_μ in the neutralino and chargino mass matrices. Full expressions for the chargino, neutralino and gluino exchange contributions are found in [13]. The dependencies of the various contributions on the CP violating phases can be neatly summarized: for the chargino contribution

$$d_f^C \sim \sin \theta_\mu , \quad (11)$$

with essentially no dependence on γ_f ; whilst for the gluino contribution

$$d_f^G \sim \overline{m}_f \sin \gamma_f , \quad (12)$$

independent of θ_μ , and the neutralino contribution has pieces that depend on both $\sin \theta_\mu$ and $\overline{m}_f \sin \gamma_f$. All three contributions can be important (including the neutralino contribution, in the case of the electron EDM), and depending on $\sin \theta_\mu$ and $\sin \gamma_f$, they can come in with either the same or opposite signs. In particular, $\text{sign}[d_f^C/d_f^G] = \text{sign}[\sin \theta_\mu / \sin \gamma_f]$. For the mass ranges we consider, the dipole moments fall as the sfermion masses are increased, and sfermion masses in the TeV range can bring these contributions to the neutron and electron electric dipole moments below the experimental bounds of $|d_n| < 1.1 \times 10^{-25} e \text{ cm}$ [23] and $|d_e| < 1.9 \times 10^{-26} e \text{ cm}$ [24], even for large values of the CP violating phases [13]. However, these large sfermion masses are inconsistent with the cosmological bounds mentioned above, where sfermion masses must be relatively close to the bino mass in order to keep the relic density in check.

We proceed as follows. We fix the value of γ_b and take $|\mu| = 3000 \text{ GeV}$. Then for several values of \overline{m}_t between 0 and 1500 GeV, we determine the upper bound on $m_{\tilde{B}}$, as a function of θ_μ . As we vary θ_μ across its full range, \overline{m}_b and γ_t change, and this affects the annihilation rate (see (9) and following discussion), and consequently the bound on $m_{\tilde{B}}$. Taking $m_{\tilde{B}}$ at its maximum value allows us to take M_L^2 as large as possible; although the electric dipole moments depend on $m_{\tilde{B}}$ as well, we find that the dependence on M_L^2 is sufficiently strong that the EDM's take their minimum values for the maximum values of $m_{\tilde{B}}$ and M_L^2 . We then compute the quark and electron EDM's as a function of θ_μ and \overline{m}_t , and use the nonrelativistic quark model to relate the neutron EDM to the up and down-quark EDM's via

$$d_n = \frac{1}{3}(4d_d - d_u). \quad (13)$$

If we find no region of the $\theta_\mu - \bar{m}_t$ parameter space which satisfies both the neutron and electron EDM bounds, we decrease γ_b and repeat the procedure. In practice, we find the bound on the neutron EDM the more difficult of the two to satisfy, and every region of the parameter space we show which produces an acceptable neutron EDM also produces a sufficiently small electron EDM. We will consequently drop further discussion of the electron EDM and concentrate on the neutron.

For the large value of $|\mu| = 3000$ GeV, we find that the largest contribution to the neutron EDM comes either from gluino exchange (for the more negative values of θ_μ) or chargino exchange (for the more positive values of θ_μ), and that the value for $|d_n|$ is too large unless γ_b takes a relatively small value. In particular, we find non-negligible experimentally acceptable regions of the parameter space only for $\gamma_b \lesssim \pi/25$. In Figure (3) we show a contour plot of the neutron EDM as a function of θ_μ and \bar{m}_t for $\gamma_b = \pi/40$. The shaded regions demarcate the range of θ_μ for this choice of γ_b . Much of this range produces a sufficiently small $|d_n|$. As we increase \bar{m}_t , the \tilde{d} and \tilde{u} masses become large and $|d_n|$ falls. As we move to values of \bar{m}_t greater than ~ 1500 GeV, we begin to require a significant tuning of M_L^2 to produce $\Omega_{\tilde{B}} h^2 < 1/4$.

Near the boundaries of the allowed range of θ_μ , γ_t approaches $\gamma_b \pm \pi/2$. As we explain above, the top contribution to the bino annihilation drops off as we move away from $\gamma_t = \pi$, and the upper bound on the bino mass, and sfermion masses, falls. It is these lower values for the sfermion masses which are primarily responsible for the sharp rise in $|d_n|$ near the boundaries. There is one last subtlety which requires mention. For every set of the parameters $\{|\mu|, \bar{m}_t, \gamma_b, \theta_\mu\}$, there are two possible values for $\{\gamma_t, \bar{m}_b\}$. We find that one of the two sets of values always gives a smaller $|d_n|$, given the cosmological constraints on the bino and sfermion masses, and it is these smaller values of $|d_n|$ which we plot.

We repeat this procedure for $|\mu| = 1000$ GeV. For lower $|\mu|$, the chargino exchange contribution is enhanced relative to the gluino exchange contribution. In Figure (4), we show a contour plot of d_n for $\gamma_b = \pi/8$. At $(\bar{m}_t = 1500 \text{ GeV}, \theta_\mu = 0)$, d_n vanishes, as $\bar{m}_b = \theta_\mu = \sin \gamma_t = 0$. Of course the gluonic and quark color dipole contributions to d_n will not vanish everywhere along the contour $d_n = 0$, but their contributions are $\lesssim 3 \times 10^{-26} e \text{ cm}$ in the region plotted. Also, the gluino and chargino contributions scale differently as $m_{\tilde{B}}$ and $m_{\tilde{f}}$ are changed. If we take $m_{\tilde{B}}$ and $m_{\tilde{f}}$ less than their maximum values in Figure (2), the contours of Figure (4) will shift, by $\lesssim 1. \times 10^{-25}$. One should therefore concentrate on the qualitative features of Figure (4), as the exact positions of the contours are not significant.

In summary, we have found that CP violating phases in the MSSM can significantly affect the cosmological upper bound on the mass of an LSP bino. In particular, taking the

maximal value $\pi/2$ for the phase γ_b of the off-diagonal component of the $T_3 = -1/2$ sfermion mass matrices pushes the upper bound on $m_{\tilde{B}}$ up from ~ 250 GeV to ~ 650 GeV. When we additionally consider constraints on neutron and electron electric dipole moments, we find the upper bound on $m_{\tilde{B}}$ is reduced to ~ 350 GeV. Various combinations of the CP violating phases are constrained as well: $|\theta_\mu| \lesssim 0.3$ and $|\gamma_b| \lesssim \pi/6$ for $|\mu| \gtrsim 1000$ GeV, while γ_t and θ_A are essentially unconstrained. We note that although the bounds on θ_μ and γ_b are small, they are much larger than the values of order 10^{-3} typically considered.

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References

- [1] S.S. Gershtein and Ya.B. Zeldovich, JETP Lett. **4** (1966) 120; R. Cowsik and J. McClelland, Phys. Rev. Lett. **29** (1972) 669; A.S. Szalay, G. Marx, Astron. Astrophys. **49** (1976) 437.
- [2] P. Hut, Phys. Lett. **69B** (1977) 85; B.W. Lee and S. Weinberg, Phys. Rev. Lett. **39** (1977) 165; E.W. Kolb and K.A. Olive, Phys. Rev. **D33** (1986) 1202; E: **34** (1986) 2531; L.M. Krauss, Phys. Lett. **128B** (1983) 37.
- [3] M. Srednicki, R. Watkins and K.A. Olive, Nucl. Phys **B310** (1988) 693.
- [4] T. Falk, K.A. Olive, and M. Srednicki, Phys. Lett. **B339** (1994) 248.
- [5] H. Golberg, Phys. Rev. Lett. **50** (1983) 1419; L.M. Krauss, Nucl. Phys. **B227** (1983) 556.
- [6] J. Ellis, J. Hagelin, D.V. Nanopoulos, K.A. Olive and M. Srednicki, Nucl. Phys. **B238** (1984) 453.
- [7] K.A. Olive and M. Srednicki, Phys. Lett. **B230** (1989) 78; Nucl. Phys. **B355** (1991) 208.
- [8] K. Greist, M. Kamionkowski, and M.S. Turner, Phys. Rev. **D41** (1990) 3565.

- [9] T. Falk, R. Madden, K.A. Olive, and M. Srednicki, Phys. Lett. **B318** (1993) 354.
- [10] R. Flores, K.A. Olive, and S. Rudaz, Phys. Lett. **B232** (1989) 377.
- [11] J. Ellis, S. Ferrara, and D.V. Nanopoulos, Phys. Lett. **114B** (1982) 231; W. Buchmüller and D. Wyler, Phys. Lett. **121B** (1983) 321; J. Polchinski and M. Wise, Phys. Lett. **125B** (1983) 393; F. del Aguila, M. Gavela, J. Grifols, and A. Mendez, Phys. Lett. **126B** (1983) 71; D.V. Nanopoulos and M. Srednicki, Phys. Lett. **128B** (1983) 61.
- [12] M. Dugan, B. Grinstein & L. Hall, Nucl. Phys. **B255**, 413 (1985).
- [13] Y. Kizukuri & N. Oshimo, Phys. Rev. **D45** (1992) 1806; **D46** (1992) 3025.
- [14] J. Ellis and S. Rudaz, Phys. Lett. **B128** (1983) 248.
- [15] J. McDonald, K.A. Olive, and M. Srednicki, Phys. Lett. **B283** (1992) 80; J. Lopez, D.V. Nanopoulos, and K. Yuan, Phys. Lett. **B267** (1991) 219; and Nucl. Phys. **B370** (1992) 445; M. Drees and M.M. Nojiri, Phys. Rev. **D47** (1993) 376; R.G. Roberts and L. Roszkowski, Phys. Lett. **B309** (1993) 329.
- [16] K. Greist and D. Seckel, Phys. Rev. **D43** (1991) 3191.
- [17] S. Mizuta and M. Yamaguchi, Phys. Lett. **B298** (1993) 120.
- [18] J. Alitti *et al.*, Phys. Lett. **B235** (1990) 363.
- [19] S. Weinberg, Phys. Rev. Lett. **63** (1989) 2333.
- [20] R. Arnowitt, M. Duff, and K. Stelle, Phys. Rev. **D43** (1991) 3085.
- [21] E. Braaten, C.S. Li, and T.C. Yuan, Phys. Rev. Lett. **64** (1990) 1709; J. Dai, H. Dykstra, R.G. Leigh, S. Paban, and D.A. Dicus, Phys. Lett. **B237** (1990) 216; **B242** (1990) 547(E).
- [22] R. Arnowitt, J. Lopez, & D. V. Nanopoulos, Phys. Rev. **D42**, 2423 (1990).
- [23] I. S. Altarev *et al.*, Phys. Lett **B276** (1992) 242.
- [24] K. Abdullah *et al.*, Phys. Rev. Lett. **65**, 2347 (1990).

Figure Caption

Fig. 1) Upper limits on the bino mass as a function of the off-diagonal element \overline{m}_t in the top squark mass² matrix, for various values of γ_b , the argument of the off-diagonal element of the $T_3 = -1/2$ sfermion mass² matrix. Also shown is the lower bound (lowest curve) on the bino mass assuming $\gamma_b = \pi/2$. The value of $|\mu|$ was chosen to be 3000 GeV.

Fig. 2) As in Fig. 1, with $|\mu| = 1000$ GeV.

Fig. 3) Contours of the neutron electric dipole moment, d_n , in the $\theta_\mu - \overline{m}_t$ plane in units of $10^{-25}e$ cm. The value of $|\mu|$ was chosen to be 3000 GeV. The shaded region corresponds to values of θ_μ and \overline{m}_t which are not allowed algebraically for this value of μ and γ_b .

Fig. 4) As in Fig. 3, with $|\mu| = 1000$ GeV.







